

Undersaturation of quarks at early stages of relativistic nuclear collisions: the hot glue initial scenario and its observable signatures

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The early stage of high multiplicity nuclear collisions is represented by a nearly quarkless, hot, deconfined pure gluon plasma. This new scenario should be characterized by a suppression of high p_T photons and dileptons as well as by reduced baryon to meson ratios. We present the numerical results for central Pb+Pb collisions at the LHC energies by using the ideal Bjorken hydrodynamics with time-dependent quark fugacity. It is shown that about 25% of final total entropy is generated during the hydrodynamic evolution of chemically undersaturated quark-gluon plasma.

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1 Introduction

The proper understanding of the initial and the early stage of ultra-relativistic pp-, pA- and heavy ion AA- collisions is a topic of great importance for our understanding of hot and dense QCD matter formed in the laboratory and in the early universe. One of the central questions is how the initial highly nonequilibrium system evolves to a state of partial thermodynamic equilibrium at later stages of nuclear collisions. There exist several models which describe the initial state in terms of non-equilibrium parton cascades (Wang & Gyulassy 1991; Xu & Greiner 2005), minijets (Eskola & Kajantie 1997), color glass condensate (McLerran & Venugopalan 1994), coherent chromofields (Magas et al. 2001; Mishustin & Kapusta 2002) etc. It is commonly believed that the strong non-equilibrium effects persist only for a short time $\sim 1/Q_s$, where $Q_s \simeq 1 - 2$ GeV is the so-called saturation scale (Gribov et al. 1983), but at later times the system reaches a state of a partial thermodynamic equilibrium.

Relatively large gluon-gluon cross sections lead to the idea (van Hove & Pokorski 1975) that the gluonic com-

ponents of colliding nucleons interact more strongly than the quark-antiquark ones. Then the two-step equilibration of the quark-gluon plasma (QGP) was proposed (Raha 1990; Shuryak 1992; Alam et al. 1994; McLerran & Venugopalan 1994; Krasnitz & Venugopalan 2001). In this approach the gluon thermalization takes place at the proper time $\tau_g < 1 fm/c$ and the (anti)quark equilibration occurs at $\tau_{th} > \tau_g$.

The surprising result, that only very few soft quarks are present at an early stage of a relativistic collision, was obtained in many transport calculations (Biró et al. 1993; Roy et al. 1997; Elliott & Rischke 2000; Blaizot et al. 2013; Uphoff et al. 2015). Observable consequences of the above two-step scenario was considered by several authors, see e.g. (Strickland 1994, Kämpfer & Pavlenko 1994, Traxler & Thoma 1996; Dutta et al. 2002; Gelis et al. 2004; Scardina et al. 2013; Liu et al. 2014; Monnai 2014; Stoecker et al. 2015; Vovchenko et al. 2015a). One immediate prediction of such a "pure glue" initial scenario is suppressed yields of hard "thermal" photons and dileptons¹. Such a suppression occurs due to the reduction of the electric charge density

¹ One should distinguish such particles from photons and Drell-Yan dileptons produced in inelastic collisions of initial partons.

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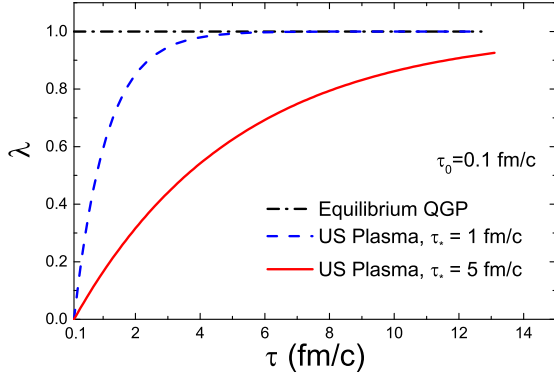


Fig. 1 (Color online) The quark fugacity λ as a function of proper time τ . The solid and dashed lines correspond to chemically undersaturated (US) plasma with parameters $\tau_* = 1$ fm/c and 5 fm/c, respectively. The dashed-dotted line corresponds to the case of chemical equilibrium ($\lambda = 1$).

as compared to chemically equilibrated quark-gluon plasma (QGP) at early stages of the reaction.

2 Evolution of undersaturated QGP in nuclear collision

Below we assume that a thermally (but not necessary chemically) equilibrated QGP is created initially in a nuclear collision. In this section we use the parameters typical for central Pb+Pb collisions at the LHC energy $\sqrt{s_{NN}} = 2.76$ TeV. By using the one-dimensional scaling hydrodynamics (Bjorken 1983) we consider the space-time evolution of QGP produced at the proper time $\tau = \tau_0$. We adopt the equation of state of an ideal gas of massless gluons, quarks and antiquarks. At zero net baryon density one may describe deviations from chemical equilibrium for quarks and antiquarks² by introducing the quark fugacity λ . Within these approximations, the following relations for the energy density ε and pressure P can be written ($\hbar = c = 1$)

$$\varepsilon = 3P = \sigma T^4, \quad \sigma = \frac{\pi^2}{30} \left(16 + \lambda \frac{21}{2} N_f \right), \quad (1)$$

where T is temperature, N_f is the number of quark flavours (unless stated otherwise, we assume that $N_f = 3$). The first and second terms in the last equality describe, respectively, the contributions of gluons and quark-antiquark pairs to the energy density. The parameter λ changes from zero for the pure gluonic system to unity for chemically equilibrated QGP.

By using the relation $P = \varepsilon/3$ and neglecting the viscosity effects one can easily get the analytic solution of the hydrodynamic equation in the Bjorken model

$$\varepsilon = \varepsilon(\tau_0) \left(\frac{\tau_0}{\tau} \right)^{4/3}, \quad (2)$$

² In accordance with the two-step approximation (see Sec. 1) we neglect deviations from chemical equilibrium for gluons, assuming that their fugacity is equal to unity during the whole process of the QGP evolution at $\tau \geq \tau_0$.

where the parameter τ_0 corresponds to the initial proper time of the hydrodynamic expansion. In principle one may determine λ and T as functions of τ by solving numerically the additional rate equation for quark density evolution (see e.g. Biró et al. 1993; Monnai 2015). The qualitative analysis can be performed by introducing the analytic parametrization (Vovchenko et al. 2015a)

$$\lambda(\tau) = 1 - \exp\left(-\frac{\tau_0 - \tau}{\tau_*}\right), \quad (3)$$

where τ_* is the model parameter characterizing the quark chemical equilibration time. Calculations of different authors lead to different estimates for τ_* , ranging from $\tau_* \sim 1$ fm/c (Ruggieri et al. 2015) to $\tau_* \sim 5$ fm/c (Xu & Greiner 2005). One should have in mind that this parameter may depend on the combination of nuclei and the bombarding energy. We expect that τ_* will be larger for peripheral events and lighter combinations of nuclei. Figure 1 shows the $\lambda(\tau)$ values for different choices of the parameter τ_* .

Introducing the quark chemical potential $\mu = T \ln \lambda$ and using thermodynamic relations, one can write down the following expression for entropy density of the QGP

$$s \simeq \frac{32\pi^2}{45} T^3 \left[1 + \lambda (0.66 - 0.16 \ln \lambda) N_f \right]. \quad (4)$$

By using (1)–(4) one can show that $s\tau$ is increasing function of τ , i.e. $s\tau \geq s_0\tau_0$, where the equality holds only in the equilibrium limit $\lambda = 1$.

Within the Bjorken model the total entropy per unit space-time rapidity $\eta = \tanh^{-1}(z/t)$ can be expressed as (Satarov et al. 2007)³

$$\frac{dS(\tau)}{d\eta} = \pi R_A^2 s(\tau) \tau, \quad (5)$$

where R_A is the geometrical radius of the colliding nuclei. Therefore, we obtain that the total entropy per unit space-time rapidity is not conserved: it gradually increases during the system expansion from the pure glue initial state. Note that this increase occurs within the ideal hydrodynamics, in absence of viscosity effects. We think that future models for extracting the viscosity values from the observed data should take into account the suppression of quarks at the initial state of a nuclear collision.

To fix the initial temperature, we assume that the Bjorken solution is valid until the freeze-out hypersurface $\tau = \tau_f$. The latter is determined by the condition $T(\tau_f) = 156$ MeV. Such a temperature has been extracted (Stachel et al. 2014) from the thermal fit of hadron ratios in the considered reaction. In our calculation we use the approximate relation (Hwa & Kajantie 1985) between the total entropy per unit space-time rapidity and the rapidity density of pions

$$\frac{dS(\tau_f)}{d\eta} = \nu \frac{dN_\pi}{dy} \Big|_{y=\eta} \simeq 1.7 \cdot 10^4, \quad (6)$$

where $\nu \simeq 6.3$ is the entropy per pion at freeze-out (Vovchenko et al. 2015a) and $dN_\pi/dy \simeq 2700$ is the observed yield of pions at midrapidity (Abbas et al. 2013).

³ We consider a purely central collision of equal nuclei.

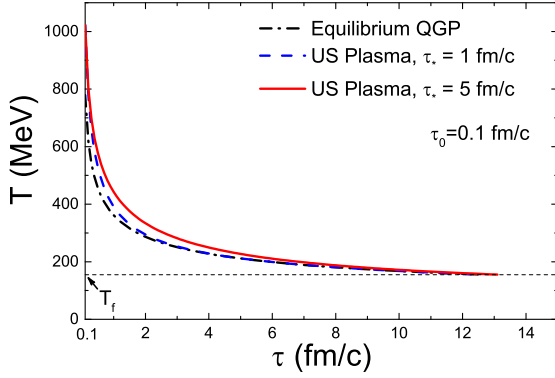


Fig. 2 (Color online) Temperature of QGP produced in central PbPb collision ($\sqrt{s_{NN}} = 2.76$ TeV) as a function of τ . The solid and dashed curves correspond to chemically undersaturated matter assuming the parameters $\tau_* = 1$ fm/c and 5 fm/c, respectively. The dashed-dotted line is calculated within the equilibrium scenario.

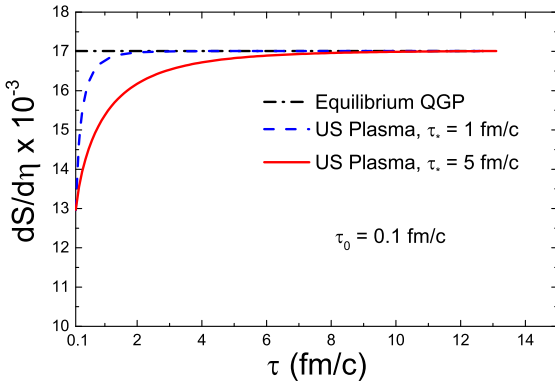


Fig. 3 (Color online) Same as Fig. 2, but for the entropy per unit space-time rapidity.

By using this procedure we have calculated the temperature and entropy density of matter in the pure glue initial scenario and compared the results with the chemical equilibrium case. Figure 2 shows the evolution of temperature for the same values of parameters τ_* and τ_0 as in Fig. 1. One can see that temperature of the US plasma is noticeably higher than in the equilibrium scenario. This increase is especially visible at $\tau \lesssim \tau_*$.

Figure 3 shows the results for the total entropy evolution. One can see that this quantity gradually increases and reaches the freeze-out value (6) during the time interval $\Delta\tau \sim \tau_*$. According to our calculations, the total increase of entropy is not sensitive to τ_* and equals about 25% of the final value.

3 Dilepton and photon spectra in the pure glue initial scenario

Hard thermal photons and dileptons are sensitive probes of hot initial stages of high energy nuclear collisions. We assume that hard dileptons are mostly produced in the decon-

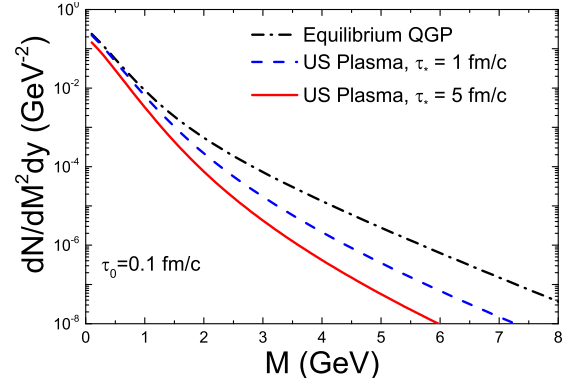


Fig. 4 (Color online) Mass distribution of thermal dileptons in central PbPb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. The solid and dashed lines are calculated with parameters $\tau_* = 1$ fm/c, and 5 fm/c, respectively. The dashed-dotted line corresponds to the chemical equilibrium scenario.

finied phase via the $q\bar{q} \rightarrow l^+l^-$ annihilation processes. Below we again apply the Bjorken hydrodynamics to describe the evolution of QGP in a heavy-ion collision. As compared to early calculations (Hwa & Kajantie 1985; Kämpfer et al. 1990) we include the additional factor λ^2 which takes into account the reduction of quark and antiquark densities in the chemically nonequilibrium case. We get the following expression for the mass spectrum of e^+e^- pairs:

$$\frac{dN_{e^+e^-}}{dM^2 dy} = \frac{\alpha^2}{\pi^2} \sum_{i=u,d,s} q_i^2 R_A^2 M \times \int_{\tau_0}^{\tau_f} \tau d\tau T(\tau) K_1[M/T(\tau)] \lambda^2(\tau), \quad (7)$$

where $\alpha = e^2$ is the electromagnetic coupling constant, q_i is the charge of the quark flavour i in units of e and $K_1(x)$ is the Macdonald function. As above, current masses of quarks are disregarded for all flavours. Note that different scenarios considered in Sec. 2 correspond to different choices of $\lambda(\tau)$ and $T(\tau)$. The results of numerical calculations are shown in Fig. 4. One can see that at $M \gtrsim 1$ GeV/ c^2 the dilepton spectra are strongly sensitive to chemical nonequilibrium effects. A special investigation shows that such dileptons are mainly produced at hot early stages of the reaction. In the pure glue initial scenario yields of hard dileptons are suppressed as compared to the equilibrium case.

To study the emission of hard thermal photons we proceed from analytic formulae for chemically equilibrated QGP suggested by Kapusta et al. (1991). In the lowest-order approximation in strong coupling constant α_s main sources of the real photon production are the qg and $\bar{q}g$ Compton scatterings as well as the $q\bar{q}$ annihilations. Attempts to consider the chemically-nonequilibrium scenario have been already made in (Strickland 1994; Kämpfer & Pavlenko 1994; Traxler & Thoma 1996; Dutta et al. 2002; Gelis et al. 2004). Following their procedure, we include additional suppression factors λ and λ^2 into the components of photon production corresponding, respectively, to the Compton scattering

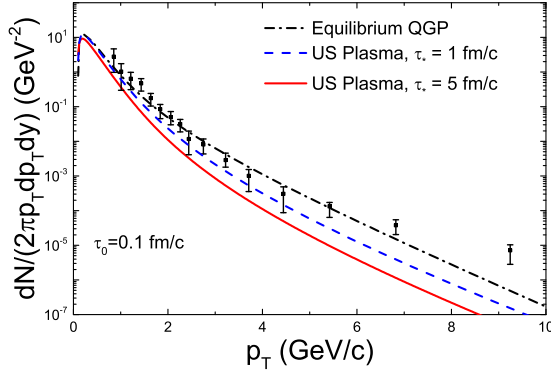


Fig. 5 (Color online) Spectrum of the thermal photons as a function of transverse momentum in central PbPb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. The solid and dashed lines correspond, respectively, to $\tau_* = 1$ fm/c and 5 fm/c. Dots show experimental data (Lohner et al. 2013) for (0–40)% most central events.

and $q\bar{q}$ annihilation terms. We have obtained the following expression for the invariant momentum distribution of hard thermal photons (Vovchenko et al. 2015b):

$$\frac{dN_\gamma}{d^2p_T dy} \approx \frac{4\alpha}{\pi^3} \sum_{i=u,d,s} q_i^2 R_A^2 \int_{\tau_0}^{\tau_f} \tau d\tau \alpha_s T^2 \times \left\{ \lambda^2 \left[\ln\left(\frac{ap_T}{\alpha_s T}\right) K_0\left(\frac{p_T}{T}\right) + \frac{bp_T}{T} K_1\left(\frac{p_T}{T}\right) \right] + \lambda \ln\left(\frac{cp_T}{\alpha_s T}\right) K_0\left(\frac{p_T}{T}\right) \right\}. \quad (8)$$

Here y and p_T are the rapidity and transverse momentum of photons (it is assumed that $p_T \gtrsim T$), and the constants $a \approx 0.20$, $b \approx 0.99$, $c \approx 0.88$. Below we disregard the temperature dependence of strong coupling constant, assuming that $\alpha_s = 0.3$.

Figure 5 shows the photon spectra calculated for the same reaction and same model parameters as above. Again one can see a noticeable suppression of high p_T photon yields as compared to the equilibrium scenario⁴. According to Fig. 5 the observed data are better reproduced for smaller values of τ_* . Note, that our calculation does not include the contribution of photons from initial parton-parton collisions. This prediction should be verified in more realistic calculations, taking into account the transverse flow of deconfined matter.

4 Evolution of the pure-gluon matter

It is instructive to consider qualitatively the evolution of the idealized pure-gluon matter created in relativistic nuclear collisions⁵. According to the QCD lattice calculations (Celik et al. 1983a, 1983b; Karsch 2002; Borsányi et al. 2012, Francis et al. 2015) this quarkless matter should

⁴ This contradicts the conclusion of Gelis et al. (2004) that chemical nonequilibrium effects do not modify significantly the photon spectra.

⁵ Such a case roughly corresponds to small rates of $gg \rightarrow q\bar{q}$ reactions. This can be simulated by choosing large τ_* within the approach developed in Sec. 2. However, we do not use now the ideal gas approximation for the gluonic matter.

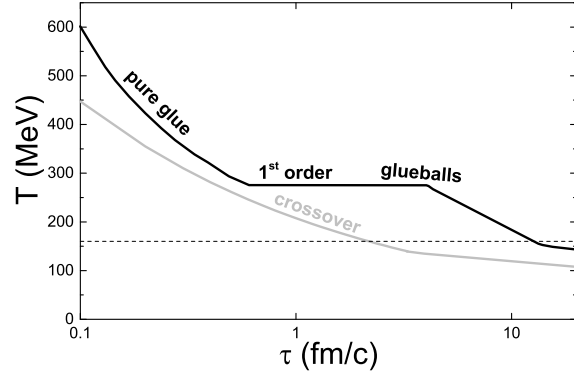


Fig. 6 Schematic picture of the temperature evolution of a high-energy collision in the pure gluon scenario with the Yang-Mills first order phase transition to glueballs.

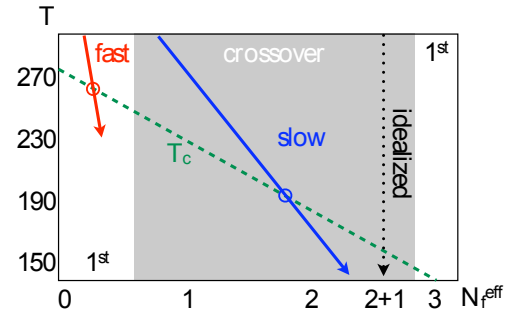


Fig. 7 (Color online) Transition temperature of baryon-free QGP versus the effective number of quark flavours.

undergo the first order phase transition at the critical temperature $T_c \approx 270$ MeV. At this temperature the deconfined pure gluon matter transforms into the confined state of the pure Yang-Mills theory, namely into a glueball fluid.

Let us assume now that a hot thermalized gluon fluid (with no quarks and antiquarks) is created at the first stage of a relativistic collision. As the system cools and expands, it may reach the mixed phase domain at $T = T_c$ and only after the gluon plasma has completely transformed into the glueball fluid, the system cools down further. The system evolution in this pure SU(3) scenario is sketched in Fig. 6. The possible appearance of super-cooled states and spinodal instabilities, associated with the first-order phase transition, can also be of particular interest.

The heavy glueballs, produced in hadronization of a pure gluon plasma, will evolve (presumably via cascade of two-body decays) into lighter states. Finally the system should decay into hadronic resonances and light hadrons. It was shown within the Frautschi - Hagedorn approach (Beitel et al. 2014) that the resulting yields of light hadrons and slopes of their spectra agree well with experimental data on heavy-ion collisions at RHIC and LHC energies.

In a more realistic scenario one should take into account that some quarks would already be produced before and during the first-order phase transition. Such a scenario could be modeled by introducing the time-dependent effective number of quark degrees of freedom. With increasing number of quark degrees of freedom the temperature of the phase transition will decrease. At some point the first-order phase transition becomes a smooth crossover. This is schematically shown in Fig. 7. The qualitative difference between the system evolution in the scenarios with the first-order and crossover transitions is demonstrated in Fig. 6. We think that the realization of a particular scenario depends on the energy and size of colliding objects. Thus, future studies of system-size and beam energy dependence of observables in nuclear collisions would be very useful.

5 Conclusions

The early stage of high multiplicity pp, pA and AA collision events can represent a new state of deconfined matter: a nearly quarkless, pure gluon plasma. According to the pure Yang - Mills lattice gauge theory, this matter undergoes, at a high temperature $T_c \approx 270$ MeV, the first-order phase transition into a confined Hagedorn-gluon fluid. Formation of such matter should lead to suppression of high p_T photons and dileptons, to reduced baryon to meson ratios, as well as to enhanced yields of heavy (e.g. charmed) hadrons. We propose to search for signatures of pure glue states in the LHC/RHIC and cosmic rays experiments.

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